

# Magnetic Monopoles as the Highest Energy Cosmic Ray Primaries

Thomas W. Kephart and Thomas J. Weiler

*Department of Physics & Astronomy, Vanderbilt University,  
Nashville, TN 37235*

## Abstract

We suggest that the highest energy  $\gtrsim 10^{20}eV$  cosmic ray primaries may be relativistic magnetic monopoles. Motivations for this hypothesis are that conventional primaries are problematic, while monopoles are naturally accelerated to  $E \sim 10^{20}eV$  by galactic magnetic fields. By matching the cosmic monopole production mechanism to the observed highest energy cosmic ray flux we estimate the monopole mass to be  $\sim 10^{10\pm 1}GeV$ .

Since their identification more than eighty years ago[1], cosmic rays have been a constant source of mystery and discovery. Of particular interest is the recent intriguing discovery of cosmic rays with energies above the GZK[2] cut-off at  $\sim 5 \times 10^{19} \text{eV}$ . Any proton energy above the cut-off is degraded by resonant scattering of the proton primary off the  $3K$  cosmic background radiation to produce the  $\Delta^*$  which then decays to nucleon plus pion. The mean free path for this process is  $\sim 6 \text{Mpc}$  for protons above the cut-off energy, and so if protons are the primaries for the highest energy cosmic rays they must either come from a rather nearby source ( $\lesssim 50 \text{Mpc}$  according to [3] and  $\lesssim 100 \text{Mpc}$  according to [4]) or have an initial energy far above  $10^{20} \text{eV}$ . Neither possibility seems likely; a  $\sim 10^{20} \text{eV}$  proton traverses a nearly straight line through the galactic magnetic field and yet no compelling local sources have been identified<sup>1</sup> near the direction of the incoming primaries[3, 4], and astrophysical mechanisms to accelerate protons to greater than  $10^{17} \text{eV}$ , let alone  $\gg 10^{20} \text{eV}$ , are speculative[6, 4]. Moreover, if  $E \gg 10^{20} \text{eV}$  protons were being emitted from a cosmically distant source, then one would also expect an accompanying flux below the GZK cut-off from roughly the same direction; this latter flux is not observed. Finally, it may be worth mentioning that a proton-induced air shower Monte Carlo[7] does not fit the shower development observed in the  $3 \times 10^{20} \text{eV}$  Fly's Eye event too well. This observation is mitigated somewhat by the fact that fluctuations in shower development are known to be large[6]. A primary heavy nucleus more closely fits the shower development of the Fly's Eye event[7]. However, a nucleus as primary has additional problems: above  $\sim 10^{19} \text{eV}$  nuclei should be photo-dissociated by the  $3K$  photon background[8] (as the nuclear lab frame energy is then above the nuclear binding energy of  $\sim 7 \text{MeV}$  per nucleon), and possibly disintegrated by the particle density ambient at the astrophysical source. Furthermore, the Fly's Eye collaboration has presented evidence that above  $\sim 10^{18} \text{eV}$  the primary composition is increasingly protons and decreasingly heavy nuclei[9].

There are now 8 events in this highest energy category, found by the AGASA[10], Fly's

---

<sup>1</sup> Two recent preprints[5] have noted that the incident directions of two of the three highest energy cosmic ray showers coincide roughly with the directions of known gamma-ray bursters (GRBs). The two directionally-coincident GRBs preceded the air showers by 5.5 and 11 months. For this association to be dynamical rather than coincidental, these two source GRBs would have to be nearby ( $\lesssim 50$  to  $100 \text{Mpc}$ ), and would have to partition nearly equal energies into gamma-rays and into  $E \gtrsim 10^{20} \text{eV}$  cosmic rays.

Eye[11], Haverah Park[12] and Yakutsk[13] collaborations. These cosmic ray detection efforts are ongoing. Furthermore, the “Auger Project” has been formed to coordinate an international effort to instrument a  $5,000 \text{ km}^2$  detector. This detector will collect five thousand events above  $10^{19} \text{ eV}$  per year [14]. Thus, there are good prospects for more cosmic ray data at these highest energies.

Another possible primary for these highest energy events is a gamma-ray. However, the time-development of the Fly’s Eye event appears inconsistent with a gamma-ray primary, as the gamma-induced shower peaks too late in the atmosphere[7]. Moreover, the mean free path for a  $\sim 10^{20} \text{ eV}$  photon to annihilate on the radio background to  $e^+e^-$  is believed to be 10 to 40  $Mpc$ [7]. It should also be noted that the density profile of the Yakutsk event[13] showed a large number of muons, which argues against gamma-ray initiation. Finally, the assignment of a neutrino as the primary is also problematic, in that the Fly’s Eye event occurs high in the atmosphere, whereas the expected event rate for early development of the neutrino-induced air shower is down from that of an electromagnetic or hadronic interaction by six orders of magnitude[7]. The acceleration problem also pertains to  $\gamma$  and  $\nu$  primaries, since  $\gamma$ ’s and  $\nu$ ’s at these energies are believed to originate in decay of  $\gtrsim 10^{20} \text{ eV}$  pions.

Given the problems with interpreting the highest energy cosmic ray primaries as protons, nuclei, photons, or neutrinos, it is not unreasonable to consider other options. Here we rekindle the idea[15] that the primary particles of the highest energy cosmic rays may be magnetic monopoles. We will show that a monopole with mass  $M \sim 10^{10} \text{ GeV}$  explains the highest energy cosmic ray data, and avoids any obvious conflict with terrestrial or astrophysical bounds.

A large motivation for this monopole hypothesis is the ease with which kinetic energies of the desired magnitude are imparted to the monopoles by cosmic magnetic fields. As pointed out by Dirac, the minimum charge for a monopole is fixed by the requirements of gauge invariance and single-valuedness of the wave function. The minimum monopole charge is  $q_M = e/2\alpha$  (which implies  $\alpha_M = 1/4\alpha$ ). In the local interstellar medium, the magnetic field  $B$  is approximately  $3 \times 10^{-6} \text{ gauss}$  with a coherence length of  $\sim 300 \text{ pc}$ [16]. Thus, a galactic monopole will typically have kinetic energy at or above

$$KE \gtrsim q_M BL \simeq 6 \times 10^{19} \text{ eV} (B/3 \times 10^{-6} \text{ gauss})(L/300 \text{ pc}).$$

Another acceleration mechanism of the right order of magnitude is provided to a monopole escaping the surface of a neutron star [17]. A monopole at the neutron star’s surface acquires a kinetic energy

$$KE = q_M BL \simeq 2 \times 10^{21} eV (B/10^{12} gauss)(L/km).$$

One might imagine that some monopoles which were initially gravitationally bound to supernova progenitor stars are ejected along the neutron star’s  $10^{12} gauss$  field lines when the star goes supernova; or that monopoles slowly migrate along the interior magnetic field of neutron stars, eventually reaching the surface where they are accelerated and ejected by the external magnetic field. (However, once monopoles have traversed a few coherence lengths in the galactic magnetic field, their energy would be expected to evolve toward the typical  $\sim 10^{20} eV$  galactic value, regardless of their initial escape velocity at the neutron star.)

We see that in both the galactic field and neutron star acceleration scenarios, there seems to be ample field strengths and field correlations to accelerate monopoles to  $\gtrsim 10^{20} eV$  energies. Furthermore, it is easy to show that radiative losses due to linear acceleration are completely negligible in the galactic acceleration scenario, and unimportant in the neutron star scenario. Thus, the “acceleration problem” for  $E \gtrsim 10^{20} eV$  primaries is easily solved. Once accelerated, the monopole retains its energy in interstellar space: inverse Compton scattering of the monopole on the 3K and diffuse photon backgrounds is negligible; for  $k_{bkgd}^0 \ll M$ , the scattering cross-section is just that of classical Thomson scattering, valid even for large coupling:  $\sigma_T = 8\pi\alpha_M^2/3M^2 \sim 2 \times 10^{-43} (M/10^{10} GeV)^{-2} cm^2$ . This cross-section is many orders of magnitude down from the pion photo-production cross-section from which the GZK cut-off derives.

To understand the expected monopole mass, number density, and mass density as a function of the monopole mass, it is necessary to review how and when a monopole is generated in a phase transition[16]. The topological requirement for monopole production is that a semisimple gauge group changes so that a  $U(1)$  factor becomes unbroken. If the mass or temperature scale at which the symmetry changes is  $\Lambda$ , then the monopoles appear as topological defects, with mass  $M = \alpha^{-1}\Lambda$ . For example, monopoles generated at the grand unification scale  $\Lambda_{GUT} \sim 10^{15} GeV$  (as determined by the running of low energy coupling constants, or by consistency with proton stability) have mass  $M \sim 10^{17} GeV$ . Such a heavy

mass remains non-relativistic after acceleration by either of the above mechanisms. Hence, a standard GUT monopole would generate no relativistic secondaries as it passes through the atmosphere, in conflict with observation. A well-known GUT example with monopole production is provided by minimal  $SU(5)$  grand unification, where a Higgs **24** breaks the  $SU(5)$  symmetry to  $SU(3) \times SU(2) \times U(1)$  at the  $\sim 10^{15} GeV$  GUT scale. On the other hand, if the symmetry breaking scale associated with the production of monopoles is below  $\sim 10^9 GeV$ , then the monopole mass is less than the monopole kinetic energy  $\sim 10^{20} eV$ , and the monopoles are relativistic. We restrict the monopole mass to  $M \lesssim 10^{11} GeV$  to ensure that the air shower induced by the monopole contains relativistic particles.

This  $M \lesssim 10^{11} GeV$  restriction also serves to avoid overclosure of the universe by an excessive monopole mass density. According to the Kibble mechanism[18], roughly one monopole is produced per horizon size at the time of the phase transition. This implies that the monopole mass density today relative to the closure value is

$$\Omega_M \sim 10^{15} (\Lambda/10^{15} GeV)^3 (M/10^{17} GeV) \quad (1)$$

(If the monopoles are relativistic, the energy density on the rhs of Eq. (1) is enhanced by mean  $\gamma_M \equiv E_M/M$ .) With the usual GUT scale of  $\Lambda \sim 10^{15} GeV$ , the fractional monopole mass density is  $\Omega_M \sim 10^{15}$ , which overcloses the universe by fifteen orders of magnitude. On the other hand, nonrelativistic monopoles less massive than  $\sim 10^{13} GeV$  do not overclose the universe.

In order to lessen the monopole density resulting from GUT-scale symmetry breaking, two approaches have been advocated in the past:

- (1) inflation is invoked after the phase transition to dilute the monopole density[19]; or
- (2) the  $U(1)$  group is broken temporarily[20], which creates cosmic string defects which connect monopoles to anti-monopoles pairwise, which then annihilate.

Here we are suggesting a third means to lower the monopole density to an acceptable level:

- (3) reduce the mass scale  $\Lambda$  of the phase transition where the  $U(1)$  first appears.

Options (1) and (2) each yields a negligible population of GUT monopoles with mass  $M \sim \Lambda/\alpha \sim 10^{17} GeV$ , which remain non-relativistic even after acceleration by the  $3 \times 10^{-6} gauss$  interstellar magnetic field. Option (3), adopted here, is more interesting. With  $\Lambda \lesssim 10^9 GeV$ ,

we are offered an abundance of relativistic monopoles well below the closure limit and yet potentially measureable, as required for our explanation of the highest energy cosmic ray events. At the end of this Letter we present as an example a simple extension of minimal  $SU(5)$  with this intermediate monopole scale<sup>2</sup>.

So that we may better assign merits and demerits to the various candidates for the highest energy primary, let us now look at the high energy cosmic ray data in some detail. Salient features of the data are:

- (i) The showers are relativistic, requiring a relativistic primary.
- (ii) There appears to be an event pile-up just below the GZK cut-off. The pile-up is apparently preceded by a dip. Around  $10^{18}eV$  both the Fly's Eye and the Akeno experiments see a spectrum that falls with energy like  $E^{-\gamma}$ , with  $\gamma = 3.2$ . At approximately  $10^{19}eV$  or just below there is an "ankle" in the data, and the slope becomes consistent with  $\gamma = 2.7$ . At around  $6 \times 10^{19}eV$  there is a cut-off, consistent with GZK.
- (iii) There appears to be a gap in the data (the statistical significance is low at present) between  $\sim 6 \times 10^{19}eV$  and the highest energy events starting above  $E \sim 10^{20}eV$ . There is a factor of three energy gap in the AGASA data between the highest energy event at  $\sim 2.2 \times 10^{20}eV$  and the second highest energy event at  $6.7 \times 10^{19}eV$ . There is a slightly larger (factor of five) gap in the Fly's Eye data,  $\sim 3 \times 10^{20}eV$  versus  $6 \times 10^{19}eV$ .
- (iv) There are 8 events above the GZK cut-off of  $E \sim 5 \times 10^{19}eV$ .
- (v) The event rate at highest energies (again, with low statistical significance) exceeds a power law extrapolation from the spectrum below the gap. The measured differential fluxes at highest energies are  $dF/dE = 5 \times 10^{-40 \pm 0.85} / cm^2 / s / sr / eV$  [10] and  $2 \times 10^{-36} / cm^2 / s / sr / eV$  [13] at  $\sim 2 \times 10^{20}eV$ , and  $7 \times 10^{-41} / cm^2 / s / sr / eV$  at  $\sim 3 \times 10^{20}eV$  [11]. The latter two flux values we obtained from [7]. We will use the range  $F_{Exp} \sim 10^{-38 \pm 2} / cm^2 / sec / sr / eV$  in what follows.
- (vi) The Fly's Eye event at  $3 \times 10^{20}eV$  comes with some detailed shower development data. For example, the peak in this air shower cascade [11] occurs at an atmospheric depth of

---

<sup>2</sup>Symmetry breaking at an intermediate scale has been invoked before in many contexts. Examples include the Peccei-Quinn solution to the strong CP problem, the right-handed neutrino scale in "see-saw" models of light left-handed neutrino mass generation, and supersymmetry breaking in a hidden sector.

$$X_{max} = 815_{-35}^{+45} + 40 \text{ g/cm}^2.$$

(vii) So far, no events are seen above the Fly’s Eye event energy at  $3 \times 10^{20} \text{ eV}$ .

Except for the highest energy cosmic ray events, the spectrum is well fit[10] by a diffuse population of protons distributed isotropically in the universe. The pion photo-production mechanism of GZK even explains the apparent pile-up[21] of events between  $\sim 10^{19} \text{ eV}$  and  $6 \times 10^{19} \text{ eV}$ . For the events above  $10^{20} \text{ eV}$ , a different origin seems to be required. That the galactic and neutron-star magnetic fields naturally impart  $10^{20}$  to  $10^{21} \text{ eV}$  of kinetic energy to the monopole, and that there appears to be an absence of events above and just below this energy, we find very suggestive. We further point out that if the monopole’s relativistic  $\gamma$ -factor is less than 10, the monopole will forward scatter atmospheric particles to  $\gamma$ -factors less than 100, insufficient for shower development. Consequently, there is an effective energy threshold for monopole-induced air showers at  $E \sim 10M$ . Thus, an apparent threshold at  $E \sim 10^{20} \text{ eV}$  may also be explained if the monopole mass is  $\sim 10^{10} \text{ GeV}$ .

Any proposed primary candidate must be able to reproduce the observed shower evolution of the  $3 \times 10^{20} \text{ eV}$  Fly’s Eye event. Protons of energy  $3 \times 10^{20} \text{ eV}$  would peak on average at  $X_{max}^p = 900 \text{ g/cm}^2$ [7], but with large fluctuations. This is later than, but marginally consistent with, the observed value of  $815 \pm 55 \text{ g/cm}^2$ . If the primary is a heavy nucleus, the average shower maximum is shifted to[7]

$$X_{max} \simeq X_{max}^p - 55(\log_{10} A) \text{ g/cm}^2$$

where  $A$  is the atomic number of the nucleus. The peak is best fit when  $A = 35$ , but such a heavy nucleus must have originated locally which is unlikely.

Does a monopole-induced air shower fit the Fly’s Eye profile? We do not know, as more theoretical work is required before this question can be answered. For a relativistic monopole primary, the electromagnetic showering property is straightforward. A magnetic monopole has a rest-frame magnetic field  $B_{RF} = q_M \hat{r}/r^2$ . When boosted to a velocity  $\vec{v}$ , an electric field  $\vec{E}_M = \gamma \beta \vec{v} \times \vec{B}_{RF}$  is generated, leading to a “dual Lorentz” force acting on the charged constituents of air atoms. Comparing this transverse  $\vec{E}$ -field to that of a relativistic particle of charge  $Ze$ ,  $E_{\perp}^{Ze} = \gamma E_{\perp, RF}^{Ze}$ , one sees that the electromagnetic energy loss of a relativistic monopole traveling through matter is very similar to the electromagnetic energy

loss of a heavy nucleus with similar  $\gamma$ -factor and charge  $Z = q_M/e = 1/2\alpha = 137/2$  (except at energies below the minimum ionizing energy, where the energy loss of slow-moving monopoles is negligible). The result[22] is a  $\sim 6 \text{ GeV}/(g \text{ cm}^{-2})$  “minimum-ionizing monopole” energy loss<sup>3</sup>. For zenith angle  $\theta_z \lesssim 60^\circ$ , the slant depth is  $(1030/\cos\theta_z) \text{ g/cm}^2$ . Integrated through the atmosphere, the total electromagnetic energy loss is therefore  $\sim (6.2/\cos\theta_z) \text{ TeV}$ . For a horizontal shower the slant depth is  $40,000 \text{ g/cm}^2$ , and the integrated energy loss is  $\sim 240 \text{ TeV}$ .

The hadronic component of the monopole shower is likely to be complicated. The monopole itself is not absorbed or destroyed in the showering process; the monopole’s integrity is guaranteed by its topology. Moreover, the monopole mass greatly exceeds the masses of the target air atoms. Thus the monopole will continuously “initiate” the shower as it propagates through the atmosphere, in contrast to the fate of the primaries in proton, nucleus,  $\gamma$ , or  $\nu$  initiated showers. For this reason, we refer to the monopole shower as “monopole-induced” rather than “monopole-initiated.” A number of unusual monopole–nucleus interactions can take place:

- (1) The interior of the monopole is symmetric vacuum, in which all the fermion, Yang–Mills, and Higgs fields of the grand unified theory coexist. Thus, even though the Compton size of the monopole is incredibly tiny, its strong interaction size is the usual confinement radius of  $\sim 1 \text{ fm}$ , and its strong interaction cross-section is indeed strong,  $\sim 10^{-26} \text{ cm}^2$  and possibly growing with energy like other hadronic cross-sections. (Multiplying this ten millibarn cross-section with the nucleon number column density,  $(1030/\cos\theta_z)(\text{g/cm}^2)/Am_N \sim 10^{26}/\cos\theta_z \text{ cm}^{-2}$ , gives the inverse mean free path for a given monopole cosmic ray to interact strongly in air. Here,  $m_N$  is the nucleon mass, and  $A = 14.5$  is the average nucleon number for air nuclei.)
- (2) S-wave scattering of monopoles is enhanced[23], leading to monopole-catalyzed baryon-violating processes with a cross-section calculated to be  $\sim 10^{-27} \text{ cm}^2$ .

---

<sup>3</sup>We note that much of the scattering cascade occurs at very high energy ( $\sqrt{s} \sim 10^4 \text{ GeV}$  for electron ionization) so in a detailed analysis we would renormalize  $\alpha$  to the value appropriate for this scale. On the other hand,  $\alpha_M$  is renormalized so as to maintain the Dirac quantization condition  $\alpha\alpha_M = 1/4$ . We also note that the scale  $s$  differs for monopole–electron and monopole–nucleon scattering by a ratio of  $m_N/m_e \sim 2000$ , which implies that the effective  $\alpha_M$  will differ slightly for these two electromagnetic processes.



(3) Besides monopole catalyzed proton and neutron decay, the relativistic monopoles considered here can also catalyze the crossed endothermic process  $e^- + M \rightarrow M + \pi + (\bar{p} \text{ or } \bar{n})$ , after which the antibaryon initiates a hadronic shower.

(4) The monopole interaction with nuclear dipole moments can cause binding of one or more nucleons by the monopole[24, 25]. If these nucleons were bound to the monopole before it was accelerated, it is likely they remain bound throughout the acceleration process. When the monopole–nucleus bound state strikes an atom, a relativistic nucleus–nucleus collisions can result.

(5) As a monopole passes through air, its interaction with the individual nucleon magnetic moments can strongly polarize the air nuclei. These deformed nuclei can then fragment[24]. For an impact parameter of  $\sim 1 \text{ fm}$ , the deformation energy for this nuclear analog of the Drell et al. process[26] is about  $30 \text{ MeV}$ .

(6) The large (azimuthal) transverse electric field of the relativistic monopole,  $E_T = \gamma e / 2\alpha r^2$ , may also polarize the air nuclei, by pushing the charged protons away from the neutrons. These polarized nuclei may then fragment.

(7) Hard coherent elastic scattering seems possible for nuclei stripped by the ionizing  $dE/dx$  process. This is most easily seen in the rest frame of the monopole where the charged nucleus will see the monopole as a magnetic bottle, spiral in toward the core of the monopole, and then be reflected by the intense gradient of the  $1/r^2$  magnetic field.

(8) A relativistic monopole needs a  $\gamma$ -factor in excess of  $4M/Am_N$  before  $M + \bar{M}$  pair production is kinematically possible. However, it may be that electroweak-scale sphaleron processes[27] could take place since the Q-value of the monopole–air nucleus interaction is  $\sim \gamma Am_N \sim \text{TeV}$ . A sphaleron has many properties of an  $M + \bar{M}$  bound state[28].

Although there has been considerable study of the interaction of nonrelativistic monopoles with matter[29], this is not so for the relativistic case. Since many energy-loss processes may be at work in monopole-induced air showers, it seems more analytic work and eventually detailed Monte Carlo studies will be required to understand air shower development. It is possible that the standard relation between the shower characteristics and the shower energy is altered. With this caution in mind, we proceed.

Using Eq.(1) and the relation  $M \sim \Lambda/\alpha$ , the general expression for the relativistic

monopole flux may be written

$$F_M = \rho_{crit} \Omega_M / 4\pi M \simeq 200 (\Lambda / 10^{15} \text{GeV})^3 / \text{cm}^2 / \text{sec} / \text{sr}. \quad (2)$$

An interesting result is obtained if we now equate this monopole flux to the measured differential flux of highest energy cosmic rays. To do so, we must assume a spectrum for the monopole flux. There is no obvious reason why monopoles accelerated by cosmic magnetic fields should have a falling spectrum, or even a broad spectrum. So we assume that the monopole spectrum is peaked in the energy half-decade 1 to  $5 \times 10^{20} \text{eV}$ . With this assumption,

$$dF_M/dE \sim F_M / 5 \times 10^{20} \text{eV} \sim 4 \times 10^{-19} (\Lambda / 10^{15} \text{GeV})^3 / \text{cm}^2 / \text{sec} / \text{sr} / \text{eV}. \quad (3)$$

Comparing this monopole flux to the measured differential flux

$(dF/dE)_{Exp} \sim 10^{-38 \pm 2} / \text{cm}^2 / \text{sec} / \text{sr} / \text{eV}$  above  $10^{20} \text{eV}$ , we find  $\Lambda \sim 3 \times 10^{8 \pm 1} \text{GeV}$ , and from this we infer  $M \sim 10^{10 \pm 1} \text{GeV}$  so the monopoles are relativistic.

The same mechanism we propose to produce the  $10^{20} \text{eV}$  monopoles, i.e., acceleration by the galactic magnetic field, will at the same time deplete the magnetic field (an inevitable consequence of energy conservation). Compatibility with the known galactic magnetic field strengths provides the ‘‘Parker bound’’ on the galactic monopole flux[30]:  $F_M^{PB} \leq 10^{-15} / \text{cm}^2 / \text{sec} / \text{sr}$ . Comparing this flux with the general monopole flux in Eq. (2), and assuming no galactic enhancement of the monopole flux, we see that the Parker bound is satisfied if  $\Lambda \leq 10^9 \text{GeV}$ , i.e. if  $M \lesssim 10^{11} \text{GeV}$ . It is very interesting that the observed flux, with the monopole hypothesis, lies just below the Parker bound. A slightly larger observed flux would violate this bound, while a slightly lower flux would not have been observed. Perhaps there exists a dynamical reason that forces the monopole flux to saturate the Parker bound.

The values of  $\Lambda$  and  $M$  inferred from Eq.(2) are obtained if the monopole density in the universe is nearly uniform. If the monopoles are concentrated in galaxies,  $\Lambda$  and  $M$  will need to be lower. If some are trapped in condensed matter (stars, etc.), or if some have annihilated, then  $\Lambda$  and  $M$  could be somewhat higher.

Other model-dependent and independent bounds exist. The monopoles considered here satisfy all the bounds discussed in Ref.[16]. There is one model-dependent bound[31] which

challenges our derived monopole flux. If galactic magnetic fields ( $\sim 10^{-6} \text{ gauss}$ ) are to grow from small seed fields ( $B_0 \sim 10^{-11}$  to  $10^{-20} \text{ gauss}$ ) via the dynamo mechanism, then the magnetic monopole flux cannot exceed

$(10^8 \text{ yr}/\tau) [(B_0/10^{-6} \text{ gauss}) + (M/10^{17} \text{ GeV})(v/10^{-3} c)^2 (l/kpc)^{-1}] \times 10^{-16} / \text{cm}^2 / \text{sec} / \text{sr}$ , where  $\tau$  is the time-scale for field generation by the dynamo,  $v$  is the monopole velocity, and  $l$  is the seed-field coherence length. Whether or not the galactic fields are derived from seed-field growth via the dynamo mechanism is uncertain, as are the values for  $\tau$ ,  $B_0$ , and  $l$ . And since approximations were made in deriving this bound (e.g. neglecting dilution of monopoles due to the universe's expansion since the era of galaxy formation, and treating  $v$ ,  $l$ , and  $\tau$  as time-independent over the era of field formation), it is not clear how binding it is. In any case, nonrelativistic monopoles may have feasted on early seed-fields to become relativistic ( $v \sim c$ ), at which point the flux bound here becomes  $10^{-16} \times (M/10^{11} \text{ GeV}) / \text{cm}^2 / \text{sec} / \text{sr}$  for the fiducial values, only slightly more restrictive than the Parker bound.

It is interesting to make a simple estimate of the monopole flux that would emanate from supernovas, should monopoles be trapped in the progenitor stars. If the mass  $M_{SN}$  of the collapsing star were composed entirely of monopoles, the number of monopoles in the stellar mass would be  $M_{SN}/M$ . If instead, the fraction of the stellar mass that is monopoles is given by  $\Omega_M$ , then we have  $\Omega_M M_{SN}/M \sim 10^{48} \Omega_M (M/10^{10} \text{ GeV})^{-1}$  as the typical number of monopoles in the supernova progenitor star, when  $M_{SN} \sim 10 M_\odot$  is taken. From observations of Sb–Sc–type spiral galaxies similar to our own, one expects a supernova in the Milky Way about every 50 years on average. If a large fraction of the bound monopoles are ejected at relativistic velocities during or after a supernova explosion and remain within the galaxy, then the mean galactic flux today is

$$F_{SN} = \Omega_M \frac{M_{SN}}{M} \frac{\tau_{gal}}{50 \text{ yrs}} \frac{c/4\pi}{V_{gal}} \sim 10^{-18} (M/10^{10} \text{ GeV})^3 / \text{cm}^2 / \text{sec} / \text{sr}, \quad (4)$$

which is again very consistent with the observed flux if the monopole spectrum peaks at  $\sim 10^{20}$  to  $10^{21} \text{ eV}$  and  $M \sim 10^{10} \text{ GeV}$ . In the last expression,  $V_{gal} \sim 4\pi R_{gal}^3/3$  is the volume of the galaxy and we take the galactic radius to be  $\sim 30 \text{ kpc}$ , and  $\tau_{gal}$  is the duration of supernova formation which we take to be  $\sim 10^{10}$  years. We have also used Eq. (1) with  $M \sim \alpha^{-1} \Lambda$  to relate the monopole fraction to the monopole mass. The use of Eq. (1) on

the  $M_{SN}$  mass scale here, or any other mass scale, is motivated by the Equivalence Principle argument that gravity does not separate matter according to its quantum numbers.

We now show by explicit example that it is easy to imagine simple GUT models where the monopoles first appear at a cosmic temperature far below the initial GUT scale breaking, with mass  $M$  therefore also far below the initial GUT scale. The field theory requirement is that a  $U(1)$  factor first appears far below the initial GUT scale breaking. Consider an extension of minimal  $SU(5)$  containing a Higgs **10** in the spectrum, in addition to the usual Higgs **24** and **5**. If the **10** gets a vacuum expectation value (VEV) first (i.e., at a high energy/temperature scale), the symmetry breaking pattern is  $SU(5) \rightarrow SU(3) \times SU(2)$ . At a lower energy scale, the standard Higgs **24** gets a VEV and the mixed quartic term  $\mathbf{10} \overline{\mathbf{10}} \mathbf{24}^2$  in the Higgs potential acts as a positive mass term for the **10** if the quartic coupling is positive, thereby driving the VEV of the **10** toward zero[32]. At some temperature  $T^*$ , the  $\langle \mathbf{10} \rangle$  VEV returns to zero. This phenomenon of “vacuum switching”[32] restores the  $U(1)_Y$  symmetry, enlarging the full vacuum symmetry to  $SU(3) \times SU(2) \times U(1)_Y$ . Finally, a Higgs **5** gets a VEV at the electroweak scale in the standard fashion to break  $SU(2) \times U(1)_Y \rightarrow U(1)_{EW}$ . Since the  $U(1)_Y$  first appears as unbroken at the scale  $\Lambda \equiv T^*$ , this marks the onset of monopole production. The VEV of the **24** sets the scale,  $\Lambda_{24}$ , of the monopole mass  $M \sim \alpha^{-1} \Lambda_{24}$  in this model <sup>4</sup>. This example constitutes a simple existence proof that light monopoles are viable field theoretically. Many more detailed and/or realistic models could be constructed.

To conclude, we have suggested that the primary particles of the highest energy cosmic rays discovered in the past several years are relativistic magnetic monopoles of mass  $M \sim 10^{10 \pm 1} GeV$ . Energies of  $\sim 10^{20} eV$  can easily be attained via acceleration in a typical galactic magnetic field; and energies even an order of magnitude higher seem typical for monopoles ejected from neutron stars. We can suggest two possible tests of this hypothesis. First of all, the distribution of galactic-field accelerated incident monopole primaries should be asymmetric and show a preference for the direction of the local galactic magnetic field. These magnetic field lines are believed to be roughly azimuthal[33]. (This also suggests that

---

<sup>4</sup>If  $\Lambda_{24}$  and  $T^*$  are very different, then the estimate of the cosmic monopole density via the simple, single-scale Kibble mechanism must be modified.

anti-monopoles and monopoles should mainly arrive from opposite hemispheres, assuming we are not located near the edge of a magnetic domain. A forward-backward asymmetry in the event rate or energy spectrum, relative to the local galactic field direction, might suggest a net excess of monopoles or anti-monopoles in our local environment.) Secondly, the characteristics of air showers induced by monopoles may carry a distinctive signature. The electromagnetic shower should develop as if the relativistic monopole carried  $\sim 137/2$  units of electric charge ( $-137/2$  for an anti-monopole). In addition, there may be several strong interaction aspects of the monopole, each contributing to monopole-induced air shower development.

## Acknowledgements

We thank Alan Barnes, Tad Pryor, Helen Quinn, Todor Stanev, David Weintraub, and the Vanderbilt University high energy experimenters, especially John Bartelt and Med Webster, for useful discussions. This work was supported in part by the U.S. Department of Energy grant no. DE-FG05-85ER40226.

## References

- [1] V. Hess, *Physik. Zeitschr.* **XIII**, 1084 (1912), and “Selected papers on Cosmic Ray Origin Theories,” ed. S. Rosen, Dover, NY (1969).
- [2] K. Greisen, *Phys. Rev. Lett.* **16**, 748 (1966); G. T. Zatsepin and V. A. Kuzmin, *Pisma Zh. Eksp. Teor. Fiz.* **4**, 114 (1966); J. L. Puget, F. W. Stecker and J. H. Bredekamp, *Ap. J.*, **205**, 638 (1976); V. S. Berezinsky and S. I. Grigoreva, *Astron. Astrophys.*, **199**, 1 (1988); S. Yoshida and M. Teshima, *Prog. Theor. Phys.*, **89**, 833 (1993).
- [3] J. W. Elbert and P. Sommers, *Astrophys. J.* **441**, 151 (1995).
- [4] G. Sigl, D. N. Schramm, and P. Bhattacharjee, *Astropart. Phys.* **2**, 401 (1994).
- [5] M. Milgrom and V. Usov, preprint astro-ph 0505009 (1995); E. Waxman, preprint astro-ph 0505085 (1995).

- [6] See, e.g., T. K. Gaisser, “Cosmic Rays and Particle Physics,” Cambridge University Pr. (1990); or A. M. Hillas, *Ann. Rev. Astron. Astrophys.* **22**, 425 (1984).
- [7] F. Halzen, R. A. Vazquez, T. Stanev, and V. P Vankov, *Astropart., Phys.*, **3**, 151 (1995).
- [8] F. W. Stecker, *Phys. Rev.* **180**, 1264 (1969).
- [9] T. K. Gaisser et al., (Fly’s Eye Collab.) *Phys. Rev.* **D47**, 1919 (1993); D. J. Bird et al., (Fly’s Eye Collab.) *Phys. Rev. Lett.* **71**, 3401 (1993).
- [10] S. Yoshida, et al., (AGASA Collab.) *Astropart., Phys.* **3**, 105 (1995); N. Hayashida et al., *Phys. Rev. Lett.* **73**, 3491 (1994).
- [11] D. J. Bird et al., (Fly’s Eye Collab.) second ref. in [?]; *Astrophys. J.* **424**, 491 (1994); *ibid.* **441**, 144 (1995).
- [12] G. Brooke et al. (Haverah Park Collab.), Proc. 19th Intl. Cosmic Ray Conf. (La Jolla) **2**, 150 (1985); reported in M. A. Lawrence, R. J. O. Reid, and A. A. Watson (Haverah Park Collab.), *J. Phys. G* **17**, 733 (1991).
- [13] N. N. Efimov et al., (Yakutsk Collab.) ICRR Symposium on Astrophysical Aspects of the Most Energetic Cosmic Rays, ed. N. Nagano and F. Takahara, World Scientific pub. (1991); and Proc. 22nd ICRC, Dublin (1991).
- [14] J. W. Cronin and A. A. Watson, announcement from the Giant Air Shower Design Group (recently renamed the “Auger Project”), October 1994.
- [15] N. A. Porter, *Nuovo Cim.* **16**, 958 (1960).
- [16] E. W. Kolb and M. S. Turner, “The Early Universe,” Addison-Wesley PUB., NY (1991).
- [17] J. A. Harvey *Nucl. Phys.* **B236**, 25 (1984).
- [18] T. W. B. Kibble, *J. Phys.* **A9**, 1387 (1976).
- [19] A. H. Guth, *Phys. Rev.* **D23**, 347 (1981).

- [20] P. Langacker and S.-Y. Pi, *Phys. Rev. Lett.* **45**, 1 (1980); T. Farris, T. W. Kephart, T.-C. Yuan, and T. J. Weiler, *Phys. Rev. Lett.* **68**, 564 (1992); V. Dixit and M. Sher, *Phys. Rev. Lett.* **68**, 560 (1992).
- [21] C. T. Hill and D. N. Schramm, *Phys. Rev.* **D31**, 564 (1985); F. A. Aharonian and J. W. Cronin, *Phys. Rev.* **D50**, 1892 (1994).
- [22] Particle Data Group, *Phys. Rev.* **D50**, 1173 (1994).
- [23] V. Rubakov, *JETP Lett.* **33**, 644 (1981); *Nucl. Phys.* **B203**, 311 (1982); C. G. Callan, Jr., *Phys. Rev.* **25**, 2141 (1982).
- [24] G. Giacomelli, in “Theory and Detection of Magnetic Monopoles in Gauge Theories”, ed. N. Craigie, World Scientific pub., 1986.
- [25] K. Olaussen, H. A. Olsen, P. Osland, and I. Overbo, *Phys. Rev. Lett.* **52**, 325 (1984);
- [26] S. D. Drell, N. M. Kroll, M. T. Mueller, S. J. Parke, and M. A. Ruderman, *Phys. Rev. Lett.* **50**, 644 (1983);
- [27] A sphaleron is the minimum-energy baryon- and lepton-number violating classical field configuration of the standard model. An overview of sphaleron physics can be found in ref. [16].
- [28] M. Hindmarsh and M. James, *Phys. Rev.* **D49**, 6109 (1994); T. Vachaspati, *Phys. Rev. Lett.* **73**, 373 (1994);
- [29] Proc. For reviews see “Magnetic Monopoles”, ed. R. A. Carrigan, Jr., and W. P. Trower, ed. Plenum Press, 1983; Proc. “Monopole ‘83”, ed. J. L. Stone, Plenum Press pub., 1984; Proc. “Theory and Detection of Magnetic Monopoles in Gauge Theories”, ed. N. Craigie, World Scientific pub., 1986.
- [30] E. N. Parker, *Astrophys. J.* **160**, 383 (1970); M. S. Turner, E. N. Parker, and T. Bogdan, *Phys. Rev.* **D26**, 1296 (1982).
- [31] F. C. Adams et al., *Phys. Rev. Lett.* **70**, 2511 (1993).

- [32] T. W. Kephart, T.-C. Yuan, and T. J. Weiler, *Nucl. Phys.* **B330**, 705 (1990), and references therein.
- [33] We thank C. Pryor for communication on this point; see, e.g., J. P. Vallee *Astrophys. J.* **366**, 450 (1991); N. Neininger, *Astron. Astrophys.* **263**, 30 (1992).